Fundamental Symmetries and Theory

W. C. Haxton[a]

[a] Institute for Nuclear Theory, Box 351550, and Department of Physics, Box 351560, University of Washington, Seattle, Washington 98195

Nuclei are powerful laboratories for studying fundamental symmetries because they filter and enhance specific interactions. I discuss four examples — hadronic parity violation, atomic electric dipole moments, precision β decay tests, and nuclear tests of neutrino masses — to illustrate some of the progress that has been made in the past few years.

1. Introduction

Let me begin by thanking Bernard Frois and the Organizing Committee for making it possible for us to enjoy both the 1998 International Nuclear Physics Conference and the lovely city of Paris. My charge today is to summarize fundamental symmetries and theory, a topic with a rich history in nuclear physics. A generation ago nuclear studies of β decay helped established the V-A nature of the weak interaction, the conserved vector current hypothesis, and other aspects of the weak interaction that became part of the standard model's experimental foundation. Much of the work carried on today is in searches for cracks in that foundation, subtle violations of low-energy symmetries that may indicate the nature of physics beyond the standard model. The field has become much too broad to cover adequately in one talk. Thus the best I can do this morning is to try to capture the flavor of the field by presenting four selected vignettes. Two of these — the discussions of hadronic parity violation and precision β decay experiments — focus primarily on the standard model. Two others — time reversal tests and massive neutrinos — look beyond the standard model to the new physics that we expect to characterize the next level of unification.

2. Nuclear Parity Violation

The hadronic weak neutral current is relatively difficult to isolate. Although the weak interaction mediates flavor-changing decays of the Λ and other strange baryons, there are no tree-level flavor-changing neutral currents in the standard model. Furthermore, weak radiative corrections to Z^0 exchange, while flavor changing, are GIM suppressed. Thus to see the hadronic weak neutral current one must study $\Delta S = 0$ interactions, which effectively limits one to NN interactions and nuclei. As the much stronger strong and electromagnetic interactions also contribute to NN interactions, the parity violation of the weak interaction must be exploited as the experimental filter [1].

At low energies the parity violating NN interaction can be modeled as a meson exchange

interaction, in analogy with similar descriptions of the strong interaction. One meson-nucleon vertex contains the weak interaction, while the second is strong, as depicted in Fig. 1.
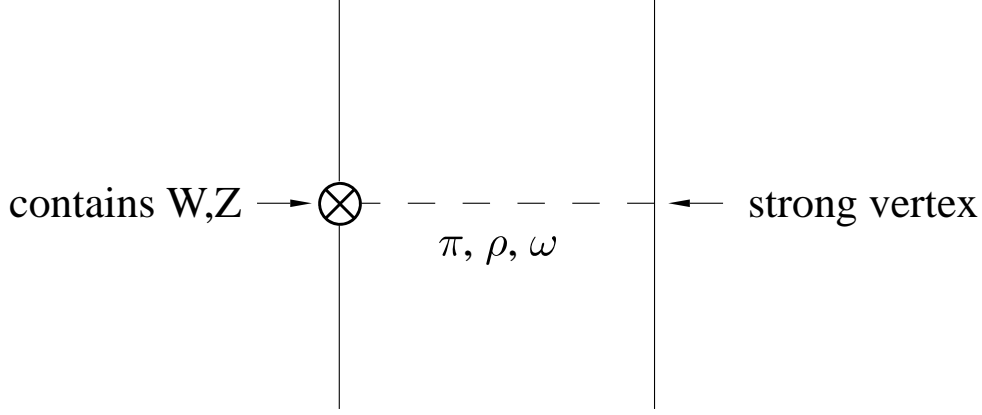

Figure 1. One-boson-exchange potential for the PNC NN interaction. One vertex is weak, containing the W or Z exchange, and one strong.

In the long-wavelength limit such a description in terms of π, ρ, and ω exchanges is sufficiently general to produce all five independent weak s-p amplitudes. To the extent that the range of the weak NN interaction is reasonably represented by such a description, it also provides a model for p-d and higher partial wave interactions.

In the standard model the low-energy hadronic weak interaction is a product of charged and neutral currents

$$L^{eff} = \frac{G}{\sqrt{2}} \left[J_W^\dagger J_W + J_Z^\dagger J_Z \right] + h.c.$$

where the charged current is comprised of $\Delta S = 0$ $\Delta I = 1$ and $\Delta S = -1$ $\Delta I = 1/2$ pieces,

$$J_W = \cos\theta_c J_W^{\Delta S=0\ \Delta I=1} + \sin\theta_c J_W^{\Delta S=-1\ \Delta I=1/2}.$$

Thus the effective Lagrangian for $\Delta S = 0$ NN interactions is

$$L_{\Delta S=0}^{eff} = \frac{G}{\sqrt{2}} \left[\cos^2\theta_c J_W^{0\dagger} J_W^0 + \sin^2\theta_c J_W^{1\dagger} J_W^1 + J_Z^\dagger J_Z \right].$$

Note that the first term above, a symmetric combination of $\Delta I = 1$ currents, has $\Delta I = 0$ and 2, only. The second term, a symmetric combination of $\Delta I = 1/2$ currents, has $\Delta I = 1$ but is suppressed by the square of the Cabibbo angle. Thus we conclude that the $\Delta I = 1$ weak NN interaction should be dominated by the third (neutral current) term above. In meson exchange models this is the channel dominated by $\pi^\pm$ exchange, the longest range component of the NN interaction. We conclude that the weak πNN coupling f_π contains the neutral current interaction we are seeking.

This coupling is the natural meeting point between experiment and nuclear theory on one hand, and the standard model on the other. As in the work of Desplanques, Donoghue, and Holstein (DDH) [2], f_π can be estimated from standard model bare quark couplings plus calculated strong interaction dressings: calculations of the latter are clearly quite difficult and model dependent, and could yield effective πNN couplings that are quite different from the underlying bare couplings. The goal is to compare the resulting f_π to experiment, thereby testing whether our theory techniques are adequate for calculating the strong interaction corrections.

As there are five elementary s-p amplitudes, the simplest experimental strategy would be to make five independent NN measurements to separate out f_π. Unfortunately only one quantity has been determined experimentally,

$$A_L(pp) = \frac{\sigma^+(\theta) - \sigma^-(\theta)}{\sigma^+(\theta) + \sigma^-(\theta)} \sim 2 \cdot 10^{-7}$$

This has meant that the needed additional constraints must be taken from nuclear experiments, where nuclear structure uncertainties can make it difficult to extract weak meson-nucleon coupling constants reliably. Fortunately there are several cases (e.g., ^{18}F and ^{19}F) where these uncertainties can be largely eliminated through ancillary experiments.

Nuclei also offer some compensating advantages: The mixing of states of definite isospin provides an isospin filter that can be exploited to isolate specific couplings, such as f_π. Furthermore, there are attractive opportunities for using nuclear degeneracies to greatly enhance the parity violating signal.

These properties are nicely illustrated by the example of ^{18}F. The level structure and the effects of the parity violation are illustrated in Fig. 2. The observable is the circular polarization of the photons emitted in the decay of the 1081 keV 0^-0 state. Because the 0^+1 1042 keV state is only 39 keV away, almost all of the parity violation is attributable to the mixing of these two states. This has the nice consequence that this observable tests the $\Delta I = 1$ part of V^{PNC}.
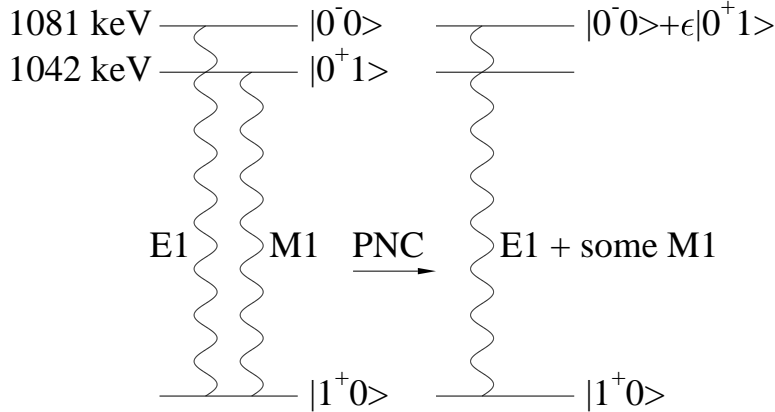
Figure 2. Schematic diagram showing the effects on parity mixing on the decay of the 1081 keV state in ^{18}F.

The enhancement comes both from the small energy denominator for the two-state mixing, and from the strength of the parity-odd M1 multipole that the parity admixture introduces into the decay of the 1081 keV state. The parity-allowed E1 decay of the 0^-0 state is very weak: the long-wavelength limit of the isoscalar E1 operator cannot generate transitions. The admixed M1 is unusually strong, on the order of 10 W.u. Thus

$$P_\gamma \sim 2Re \left[\frac{< 1^+ \parallel M1 \parallel 0^+ >}{< 1^+ \parallel E1 \parallel 0^- >} \frac{< 0^+1 \parallel V_{PNC}^{\Delta I=1} \parallel 0^-0 >}{\Delta E} \right] \sim 10^{-3}$$

The $M1/E1$ matrix element ratio is ~ 110, while $\langle V \rangle / \Delta E \sim 10^{-5}$, compared to the natural scale of weak interactions,

$$\frac{4\pi G^2 m_\pi^2}{g_{\pi NN}^2} \sim 10^{-7}.$$

It also turns out the the nuclear matrix element can be effectively measured from the axial-charge β decay of ^{18}Ne, as described in [1].

The ^{18}F result — actually an upper bound — is one of several measurements in the NN, few-body, and light nuclear systems which can be interpretted, i.e., where the nuclear physics uncertainties are sufficiently under control that the weak meson-nucleon couplings can be reliably extracted:

$$\left.\begin{array}{ll} A_L(\vec{p}+p) & 15, 45 \text{ MeV} \\ A_L(\vec{p}+^4 He) & 46 \text{ MeV} \\ A_L(\vec{p}+d) & 15 \text{ MeV (upperbound)} \\ P_\gamma(^{18}F) & \\ A_\gamma(^{19}F) & \end{array}\right\} \Rightarrow h_\rho^0 + 0.5h_\omega^0 \sim 1.2(h_\rho^0 + 0.5h_\omega^0)^{DDH}, \quad f_\pi \sim 0$$

Here f_π is the isovector weak πNN coupling and $(h_\rho^0 + 0.5h_\omega^0)$ is the combination of ρ and ω isoscalar weak couplings that arises in calculations of the ^{19}F and $\vec{p}+^4$He asymmetries. The superscript DDH denotes the best value of Ref. [2]. The net result is a surprise: although the isoscalar combination is about as predicted by DDH, f_π is consistent with zero and no larger than about one-third f_π^{DDH}.

So a summary of things as of a year ago would be:

• No evidence has been found for the neutral hadronic current.

• The isospin anomaly, $f_\pi^{PNC} \lesssim \frac{1}{3}f_\pi^{DDH}$ while the isoscalar combination has the expected size, is superficially similar to the $\Delta I = 1/2$ rule in strangeness-changing decays: the isospin dependence of the meson-nucleon couplings differs significantly from those of the underlying bare couplings. Of course, in the case of the $\Delta I = 1/2$ rule, it is an enhancement in the $\Delta I = 1/2$ amplitudes.

• Presumably the explanation for both effects must be connected with some interesting strong interaction physics occurring in the weak vertices.

A new constraint was obtained recently from a novel measurement, the first detection of a nuclear anapole moment. Atomic PNC experiments are sensitive to a variety of new physics, while also testing standard model couplings to high precision. The dominate standard model interaction is

$$A(e) - V(N).$$

That is, the leading PNC amplitude involves an axial coupling of the exchanged Z^0 to the electron and a coherent vector coupling to the nucleus. The nuclear weak charge is approximately its neutron number. About a year ago the Boulder group [3] announced an atomic PNC result of unprecedented sensitivity, a $\sim 0.3\%$ measurement in ^{133}Cs. This measurement, for example, limits the scale of new Z's to lie above 1.4 TeV — a value limited by the associated atomic theory, which is believed to be accurate to $\sim 1.0\%$.
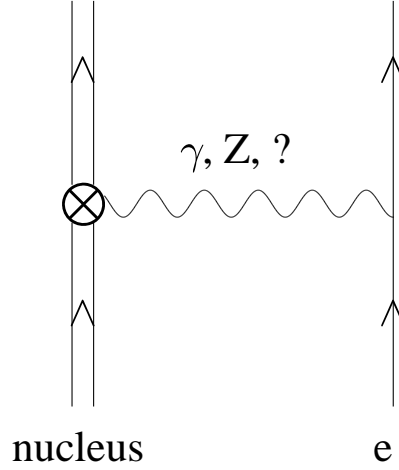

Figure 3. Atomic interactions are mediated by γ and Z^0 exchange, as well as new physics effects, denoted by ?, associated with extra Zs, leptoquarks, compositeness, etc.

The connection to nuclear PNC is the contribution associated with the axial coupling to the nucleus. There is a weak tree-level contribution to PNC from Z^0 exchange of the form

$V(e) - A(N).$

Surprisingly, even more important in a heavy nucleus are contributions due to radiative corrections. Recall that the possible static electromagnetic couplings to the nucleus are

	CJ	MJ	EJ
$J = 0$	PT		
$J = 1$	$\not{P}\,\not{T}$	PT	$\not{P}T$
$J = 2$	PT	$\not{P}\,\not{T}$	$P\,\not{T}$
$J = 3$	$\not{P}\,\not{T}$	PT	$\not{P}PT$
$\vdots$	$\vdots$	$\vdots$	$\vdots$

For example, the electromagnetic current matrix element for a spin-1/2 particle like the nucleon involves the first four entries in the above table

$$< g.s. \parallel J_\mu^{em} \parallel g.s. >= \bar{N}(p')(\overbrace{F_1\gamma_\mu}^{C0} + \overbrace{F_2\sigma_{\mu\nu}q^\nu}^{M1} + \underbrace{F_A\gamma_5\gamma_\mu + F_P\gamma_5 q_\mu}_{E1:\not{P}T anapole} + \underbrace{F_T\gamma_5\sigma_{\mu\nu}q^\nu}_{C1:\not{P}\not{T}edm})N(p)$$

where current conservation demands $F_A = \frac{q^2}{2M}F_P$. This PNC term thus can be written

$$< p' \parallel J_\mu^{em} \parallel p >_{\not{P}T} = \frac{a(q^2)}{M^2}\bar{N}(p')(\not{q}q_\mu - q^2\gamma_\mu)\gamma_5 N(p)$$

where $a(q^2)$ is the anapole moment, first discussed by Zeldovich [4, 5]. It transformations under rotations like the nucleon spin. It vanishes for on-shell photons, but for virtual photons generates a contact interaction between the scattered electron and the nucleon.

Figure 4a shows one contribution to the nucleon anapole moment: it arises as a weak radiative correction. Fig. 4b shows a similar weak radiative correction that cannot be written as a nucleon electromagnetic moment. This shows that the anapole moment is not really a measurable (i.e., gauge invariant) quantity, though for a nucleus the dominant contribution to the anapole moment is separately gauge invariant. Fig. 4c is the usual tree-level V(e)-A(N) Z^0 exchange, where the axial nuclear coupling is isovector, which combines with the electron–nuclear anapole interaction (Fig. 4d) to give the nuclear-spin-dependent parity violation in an atom.

This raises two issues: 1) how does a composite object like a nucleus generate an anapole moment? 2) how does the anapole contribution compare to the tree-level V(e)-A(N) contribution? There are three separate contributions to the nuclear anapole moment:
1) The one-body contribution:

$$< g.s.|\sum_{i=1}^{A}(a_s + a_v\tau_3(i))\vec{\sigma}(i)|g.s. >$$

is a sum over the anapole moments of the individual nucleons. As nucleons occupy spin-paired orbitals, this can be roughly thought of as the anapole moment of the last, unpaired nucleon (in the same sense that nuclear magnetic moments can be viewed in this way). If Fig. 4a is evaluated [6] for a pion loop, one finds that the isoscalar anapole moment is much larger than the isovector, $a_s \gg a_v$.
2) The exchange-current contribution, consisting of diagrams where the photon couples to a meson in flight between two nucleons, or to a nucleon-antinucleon pair.
3) The nuclear polarizability depicted in Fig. 5: the E1 anapole operator couples the unperturbed ground state to the opposite-parity components in the nuclear wave function that arise from the hadronic weak interaction. This third term easily dominates the anapole moment of heavy nuclei and is responsible for the anapole moment's $A^{2/3}$ growth, where A is the mass number.

In the Cs experiment, 7000 hours of data taking yielded the hyperfine-dependent (i.e., nuclear spin-dependent) contribution to the atomic PNC signal. The result can be expressed as the strength of the nuclear-spin-dependent electron-nucleus contact interaction

$$H_W^{e-Nuc} \equiv \frac{G_F}{\sqrt{2}}\kappa\vec{\alpha} \cdot \vec{I}\rho(r)$$

$$\kappa = \underbrace{\kappa_{V(e)-A(N)}^Z}_{0.013} + \kappa_{anapole} = \underbrace{0.112 \pm 0.016}_{7\sigma}$$

The Z^0 contribution depends on the Weinberg angle (taken to be $\sin^2\theta_W = 0.223$) and on the ^{133}Cs ground state matrix element of the Gamow-Teller operator, which was taken
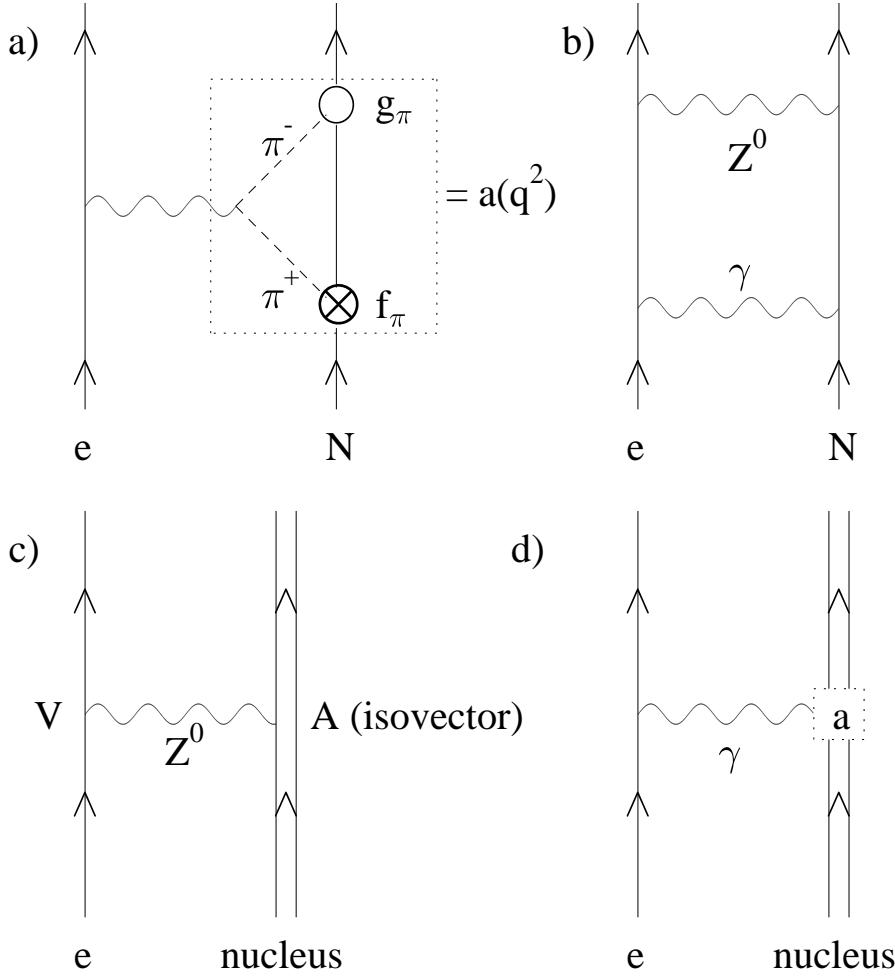
Figure 4. The nuclear anapole (a) and box (b) weak radiative corrections to the hydrogen atom. The nuclear-spin-dependent contribution to atomic PNC include the V(e)-A(nucleus) Z^0 exchange (c) and the electron interaction with the nuclear anapole moment (d).

from the shell model calculation of [6]. The extraction of κ from the Boulder result requires an atomic calculation. The above value is taken from Flambaum and Murray [7]. As the extracted κ is an order of magnitude larger that the tree-level Z^0 contribution, the anapole contribution has been clearly seen. This places the following constraint on the weak meson-nucleon coupling constants

$$f_\pi - 0.21(h_\rho^0 + 0.59h_\omega^0) = (0.99 \pm 0.16) \times 10^{-6}$$

where again a familiar combination of isoscalar weak couplings arises. (This result was calculated with the same shell model techniques as in [6], but includes the ρ and ω exchange contributions to the nuclear polarizability. Note that the average nuclear excitation energy was changed from the value used in [6], 15.2 MeV, to 9.3 MeV.) In Fig. 6 this new constraint is shown along with those from three other PNC experiments (^{18}F,
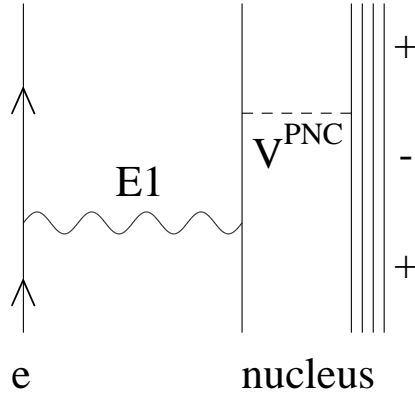
Figure 5. The PNC nuclear polarizability contribution to the electron's interaction with the nuclear anapole moment. This diagram generally dominates the anapole moment of heavy nuclei. The nuclear E1 operator is an unfamiliar one: the usual operator vanishes due to the extended Siegert's theorem (see [6]). The evaluation of the polarizability requires a summation over a complete set of intermediate nuclear states. In the fully-interacting shell model approach of [6], this sum was evaluated by closure, assuming an average excitation energy for the intermediate nuclear states. Because of properties of the E1 operator in the standard shell model space for ^{133}Cs (there are no nonzero matrix elements of J=1 odd-parity operators in the $1g_{7/2}2d_{5/2}3s_{1/2}2d_{3/2}1h_{11/2}$ space), the polarizability then reduces to an expectation value of a two-body operator.

^{19}F, $\vec{p}+^4$He), all of which test approximately the same combination of weak couplings. It is immediately apparent that the ^{133}Cs anapole moment is not in agreement with our tentative conclusion that f_π is considerably below the DDH value.

Until the ^{133}Cs anapole measurement there was very little redundancy among the experimental constraints shown in Fig. 6. We now have to grapple with an inconsistency whose origin is unknown. It could be that one of the experiments is wrong. The theory underlying Fig. 6 also has its weak points. The interpretation of the anapole moment measurement depends on the accuracy of the calculation of the ^{133}Cs polarizability. The evidence that this has been done well is the reasonable agreement between two independent calculations [6, 8], but this should be further explored. There are additional issues of concern — the use of bare operators in the nuclear calculations, strange quark contributions to the hadronic matrix elements, etc. — that require more technical discussion than is possible here.

There are several possibilities for new experiments that could greatly clarify matters. Perhaps the most important are possibilities for measuring the np weak interaction either directly or in a few-body system that can be reliable interpretted. Efforts are underway to measure the PNC $\vec{n}$ spin rotation in ^{4}He [9], and a proposal has been made to measure A_γ for $\vec{n} + p \rightarrow d + \gamma$ [10]. There is also an important experiment nearing completion at TRIUMF in which A_L for $\vec{p} + p$ is being measured at medium energies, thereby testing the combination of vector meson PNC couplings that contribute to p-d wave interference [11]. Perhaps our picture of hadronic PNC will be clearer at the time of INPC2001.
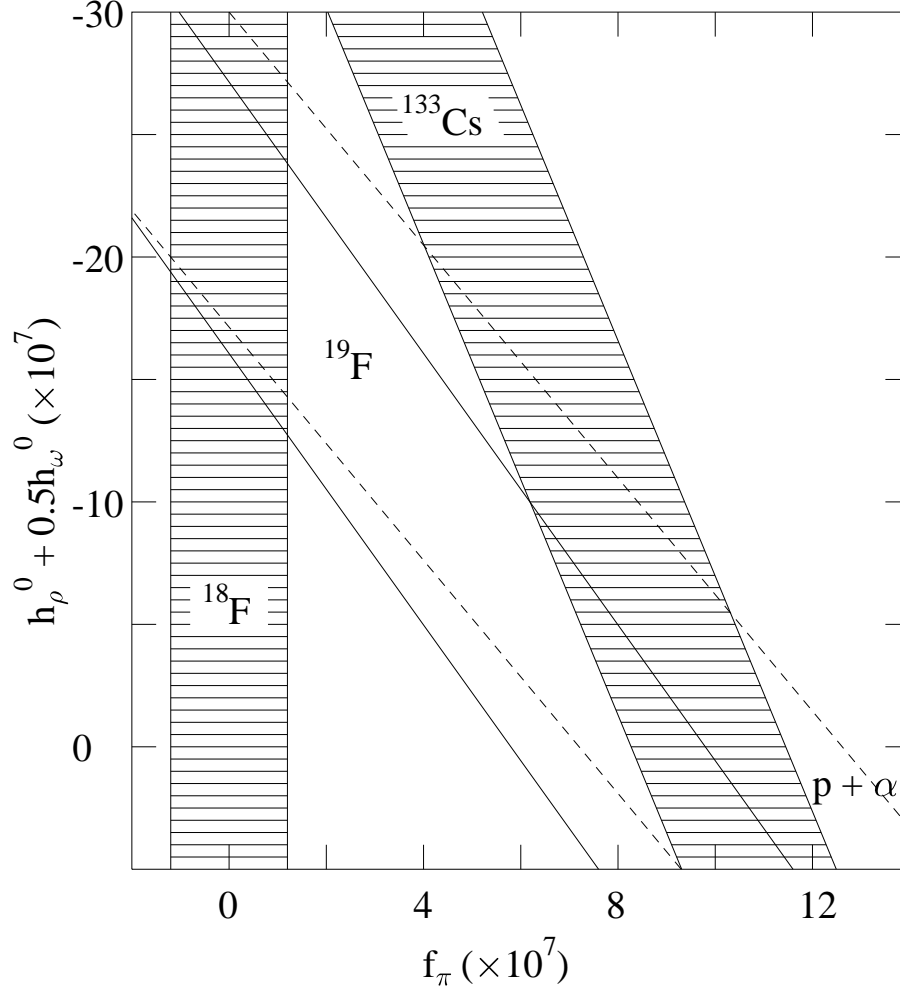
Figure 6. Experimental constraints on the isovector pion and isoscalar heavy meson weak couplings. The indicated regions correspond to the 1σ experimental errors, only; no estimate of the additional theory errors has been made.

3. Atomic Nuclei and T Violation

Recall our table of static electromagnetc couplings:

	$\underline{CJ}$	$\underline{MJ}$	$\underline{EJ}$
$J = 0$	PT		
$J = 1$	$\not{P}\,\not{T}$	PT	$\not{P}T$
$\vdots$	$\vdots$	$\vdots$	$\vdots$

The moment of interest in this section is the $C1$, the P-odd T-odd electric dipole moment. The required CP-violating, P-violating interactions occur in nature ($K_L \rightarrow \pi\pi$) and in the standard model (through the CKM phase in the quark mass matrix and the $\bar{\theta}$ term), and there are strong reasons for believing additional sources of P-violating, CP-violating interactions exist beyond the standard model. Such interactions will generate a nucleon

edm:

$$< p'|J_\mu^{em}|p >= \ldots + \bar{N}(p')\underbrace{d(q^2)}_{edm}\sigma_{\mu\nu}q^\nu\gamma_5 N(p)$$

If a neutron or neutral atom having an edm — a charge separation along the direction of spin — is placed in an external electric $\vec{E}_{ext} = E_z\hat{z}$, the interaction is

$$H = d\vec{s}\cdot\vec{E}$$

The resulting torque makes the spin precess about $\hat{z}$. Neutron measurements by the Grenoble [12] and Gatchina [13] groups and atomic measurements by the Seattle group [14] have produced wonderfully precise results

$$|d_{\text{neutron}}| \leq 8\cdot 10^{-26}\text{e cm}$$

$$|d(^{199}\text{Hg})| \leq 1.3\cdot 10^{-27}\text{e cm}$$

It is the rapid improvement in the atomic measurements that motivates my inclusion of this topic here. In many cases the most interesting sources of CP violation within the atom are those residing within the nucleus. Thus we must deal with the nuclear physics governing CP violation to understand the implications of atomic edm measurements for underlying theories of CP violation. This involves a series of steps: understanding how the underlying model generates a nucleon edm and a CP-violating NN interaction; calculating how this interaction polarizes the nucleus to produce a nuclear edm; and calculating the atomic polarization that the nuclear edm induces.
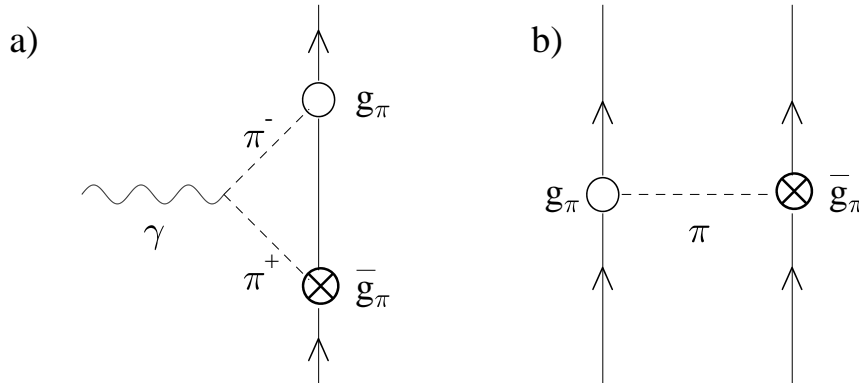

Figure 7. The nucleon edm (a) and P-violating, CP-violating NN interaction (b) that will arise from a CP-violating scalar coupling $\bar{g}$ of the pion to the nucleon.

Neutrons and atomic nuclei also provide complementary constraints on CP-violation. Consider, for example, the $\bar{\theta}$ term in the QCD Lagrangian

$$\bar{\theta}\frac{g^2}{32\pi^2}F_{\mu\nu}\tilde{F}^{\mu\nu}$$

where the neutron edm limits $|\bar\theta| \lesssim 10^{-10}$. The unnatural smallness of $\bar\theta$ is often called the strong CP problem. One low-energy consequence of $\bar\theta$ is the generation of a CP-violating scalar coupling $\bar{g}_{\pi NN}$ of the pion to the nucleon [15]

$$L_{\pi NN} \to \vec\pi \cdot \bar{N}\vec\tau(i\gamma_5 g_{\pi NN} + \bar{g}_{\pi NN})N$$

$$\bar{g}_{\pi NN} \sim 0.03\bar\theta$$

The pion coupling is important because the pion is the longest range meson and thus can produce the largest charge separation, or edm, as shown in Fig. 7a. It also should dominate the long-range CP-violating NN interaction (Fig. 7b). The neutron edm corresponding to Fig. 7a is [15]

$$d_n \sim \frac{eg_{\pi NN}\bar{g}_{\pi NN}}{4\pi^2 M} ln(\frac{M}{m_\pi})$$

where we note the chiral log. Thus an atomic edm measurement could be analyzed in terms of $\bar{g}_{\pi NN}$, which then serves as a low-energy constraint on possible models of CP violation.

The isospin dependence of the long-range πNN interaction can distinguish competing descriptions of CP-violation [16, 17]. The possibilities include isoscalar, isovector, and isotensor couplings

$$\bar{g}^0_{\pi NN}\bar{N}\vec\tau N \cdot \vec\pi$$
$$\bar{g}^1_{\pi NN}\bar{N}N\pi^0$$
$$\bar{g}^2_{\pi NN}\bar{N}(3\tau_3\pi^0 - \vec\tau \cdot \vec\pi)N$$

where $\bar\theta$ generates the first (isoscalar) coupling above. The neutron edm limit provides the constraint

$$|d_n^{exp}| \lesssim 8 \cdot 10^{-26}\text{e cm} \Rightarrow |\bar\theta| \lesssim 10^{-10}$$

A calculation shows that the ^{199}Hg edm would arise primarily from the CP-odd mixture in the ground state due to the CP-violating NN interaction, rather than from the edm of the unpaired valence nucleon. With some effort the constraint imposed by the Seattle experiment can be recast in the form

$$|d_n^{\bar\theta}| \lesssim 5 \cdot 10^{-26}\text{e cm}$$

There are several points to be made. First, the atomic and neutron edm measurements provide very similar constraints on $\bar\theta$. Second, the atomic limit is mildly dependent on the souce of the CP violation: it tightens by about a factor of four if the CKM phase is used as the source of the CP violation. Third, the isospin dependence of $\bar{g}$ leads to a distinctive scaling of the atomic result with the N and Z of the nucleus. In the case of the isoscalar coupling (i.e., $\bar\theta$), the atomic edm $\sim$ (N-Z). Thus if a nonzero neutron edm were measured, one could do additional atomic experiments looking for such a dependence, in order to clarify the origin of the neutron result.

With the recent rapid progress in atomic edm measurements, this approach now compares favorable to neutron edm experiments in sensitivity to the underlying particle

physics. But the progress is not at an end: both types of experiments continue to improve, with current efforts possibly yielding another factor of 20. In the case of the atomic experiments the detailed dependence of the resulting limits on the nuclear physics — illustrated here for $\bar{\theta}$ — points out the relevance of our field to this endeavor. The accuracy of the resulting particle physics constraints will depend on the quality of the supporting nuclear calculations. Finally, there is the exciting possibility that the enhancements we have long exploited in PNC tests might find an analog in atomic edm measurements. To date none of the nuclei identified as especially polarizable (some by enhancements factors ranging up to 10^4 [16]) have been suitable candidates for atomic measurements. But as new atomic methods are developed, this situation could certainly change.

4. β Decay Tests of Weak Interactions

As I mentioned in the introduction, β decay studies played a prominent role in establishing the experimental foundations of the standard model and, today, continue to be important as tests of new physics. There is a great deal of new activity, much of it reported in the parallel sessions of this meeting. While I no not have the time to discuss any of this in the detail it deserves, I do want to cite a few of the results of the past year that I found notable:

• Adelberger is reporting at this meeting a new study of the $e^+ - \nu_e$ correlation in $0^+ \rightarrow 0^+ \beta$ decay which tightens constraints on scalar interactions. Quoted in terms of an equivalent boson mass, the result is $M_S \gtrsim 4M_W$.

• Savard is reporting on the status of Fermi β decay tests of standard model unitarity. While there is still a modest departure from unitarity, the data for targets with $5 \lesssim Z \lesssim 26$ smoothly extrapolate to small Z. Perhaps this indicates that calculations of the weak radiative corrections and effects of isospin mixing are in reasonably good shape.

• Tritium β decay constraints on the ν_e mass continue to improve, with the Mainz and Troitsk experiments yielding constraints $\lesssim (3 - 5)$ eV. Some progress has been made in identifying effects responsible for a troubling excess of electrons very near the endpoint (the so-called negative m_ν^2 problem). Yet it is clear the spectrum shape is not yet fully understood [18].

• Atomic exchange effects — where the emitted β ray is captured into an atomic orbit while an atomic electron is kicked into the continuum — were recently measured for the first time in an allowed decay [19].

• Environmental effects in β decay — effects of the surrounding atoms of the spectrum of emitted electrons — have been seen for the first time [20]. Such effects were predicted some time ago [21].

5. Neutrino Masses

I would like to close this talk with a few comments about neutrino mass, given the excitement this field has generated these past few months. It has long been realized that neutrinos might provide a special window on physics far beyond the standard model. This became apparent from a puzzle that arose when extended models were first considered. In such a model one might hope to replace the standard model doublets with larger

multiplets in the hope of further unifiying the interactions, as depicted below

$$\begin{pmatrix} \nu \\ e \end{pmatrix} \xrightarrow[\substack{grander \\ model}]{e.g.} \begin{pmatrix} u \\ d \\ \nu \\ e \end{pmatrix}$$

The puzzle derives from the masses of the particles in the multiplet. One would expect these particles to couple to the mass-generating fields in a similar way, so that the resulting masses would be comparable up to group theory factors. But, while the electron and first generation quarks have masses on the order of an MeV, the ν_e mass is at least six orders of magnitude smaller.

Gell-Mann, Ramond, Slansky, and Yanagida [22] recognized that the neutrino is special. Unlike the other fermions which clearly have distinct antiparticles under particle-antiparticle conjugation (e.g., $e^- \to e^+$, with the electron and positron distinguished by their opposite charges), there is no obvious additive quantum number distinguishing the ν from the $\bar{\nu}$. This has the consequence that in addition to the usual Dirac mass term m_D, neutrinos can have Majorana masses that break lepton number conservation. The result is a neutrino mass matrix that, when diagonalized, yields a light neutrino mass

$$m_\nu^{light} \sim m_D \quad \underbrace{\left(\frac{m_D}{M_R} \right)}_{\text{``small parameter''}}$$

where M_R is some heavy right-handed Majorana mass characterizing scales beyond the standard model. If one fixes m_D to the charged fermion masses and m_ν^{light} to theoretical scenarios explaining the solar and atmospheric neutrino problems, then

$$M_R \sim (10^{12} - 10^{16}) \text{GeV}$$

This "seesaw" explanation of light neutrino masses suggests that neutrinos provide a window on new physics far beyond the explored low-energy world of the standard model.

To provide some picture of how these various results might fit together to form some pattern, I now discuss a recent paper by Georgi and Glashow [23]. The assumptions of their construction are:

• Three light Majorana neutrinos

• The atmospheric neutrino problem is due to $\nu_\mu \to \nu_\tau$ oscillations, since the $\nu_\mu \to \nu_e$ alternative is ruled out by the Chooz experiment.

•This oscillation is nearly maximal with $\sin 2\theta_{23} \sim 1$ and $5 \cdot 10^{-4}$ eV$^2 \lesssim \delta m_{23} \lesssim 6 \cdot 10^{-3}$ eV^2.

•The solar neutrino problem is due to oscillations with $6 \cdot 10^{-11}$ eV$^2 \lesssim \delta m^2 \lesssim 2 \cdot 10^{-5}$ eV^2.

• The neutrino masses are constrained to satisfy m$_1$+m$_2$+m$_3 \sim 6$ eV in order to generate hot dark matter for large scale structure formation (a somewhat speculative condition).

•The absence of neutrinoless double β decay requires $\langle m_\nu^{\text{Maj}} \rangle \lesssim 0.4$eV, so choose $\langle m_\nu^{\text{Maj}} \rangle \sim 0$.

•Because of the LSND/KARMEN conflict, the LSND results are not considered.

These constraints lead to a pattern of three nearly degenerate massive neutrinos with $m_i \sim M$ and a simple mass matrix that accounts for the atmospheric and solar neutrino problems through vacuum oscillations,

$$M \begin{pmatrix} 0 & \frac{1}{\sqrt{2}} & \frac{1}{\sqrt{2}} \\ \frac{1}{\sqrt{2}} & \frac{1}{2} & \frac{-1}{2} \\ \frac{1}{\sqrt{2}} & \frac{-1}{2} & \frac{1}{2} \end{pmatrix} \begin{matrix} \nu_e \\ \nu_\mu \\ \nu_\tau \end{matrix}$$

This kind of mass matrix can arise naturally in model schemes, as has been shown recently by Mohapatra and Nussinov [24]. Clearly it is just one possibility among many, but suggests that the hints of massive neutrinos we now have may yet conform to a simple pattern.

The interest in neutrino physics is likely to build in the next few years, with nuclear physics having many opportunities to contribute. The hope is that we will soon understand whether a simple pattern like that given above describes nature. So in my wishlist for our field in the years leading up to INPC2001 I would include
• LSND vs. KARMEN fully resolved
• data from SNO
• improved $\beta\beta$ decay limits reaching well beyond 0.1 eV
• some progress on very long baseline $\nu_\mu \leftrightarrow \nu_\tau$ experiments.
There is a great deal new to anticipate.

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